Plasmoid impacts on neutron stars and highest energy cosmic rays

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Particle acceleration by electrostatic polarization fields that arise in plasmas streaming across magnetic fields is discussed as a possible acceleration mechanism of highest-energy ($\gtrsim 10^{20}$ eV) cosmic rays. Specifically, plasmoids arising in planetoid impacts onto neutron star magnetospheres are considered. We find that such impacts at plausible rates may account for the observed flux and energy spectrum of the highest energy cosmic rays.

The origin of ultra-high energy cosmic rays (UHECRs), with energies up to and exceeding 10^{20} eV [1–3], remains unknown: the commonly invoked diffusive first-order Fermi acceleration of cosmic rays in a supernova shock [4–7], can accelerate particles to at most $\sim 10^{15} - 10^{16}$ eV [8]. While additional acceleration to energies ~ 100 times higher by the electric field in a pulsar-driven supernova remnant has been proposed [9], these energies are still much below the highest observed energies. Other models invoke Fermi acceleration associated with cosmological gamma ray burst sources [10] and a decay of supermassive X particles of grand unified field theories [11]. In this Letter we consider a different acceleration mechanism, based on charge polarization arising in plasmoids impacting neutron star magnetospheres.

It is well known [12,13] that an electrostatic field arises in bounded plasmas moving across the magnetic field at sub-Alfvénic velocities. The reason for this is plasma polarization caused by opposing gravitational and polarization drifts of electrons and ions that lead to the appearance of net charge near the plasma boundary. If the plasma density ρ is so high that the transverse susceptibility $\chi_\perp \equiv 4\pi\rho c^2/B^2 \gg 1$, then the electrostatic field $-\nabla\Phi = -\mathbf{V} \times \mathbf{B}/c$ where $\mathbf{V}$ is plasma flow velocity and $\mathbf{B}$ is the magnetic field; the potential drop across the plasmoid of width h in the cross-field direction (denoted by $\perp$) is $2\Phi_0 = hV_\perp B/c$ (Fig. 1).

Outside the plasmoid, the stray electrostatic field has a large component parallel to the magnetic field which causes particle acceleration along the field lines. This phenomenon has been observed in numerical simulations [14–16] which showed that charge layers can accelerate particles to relativistic energies even for relatively slow (sub-Alfvénic) plasma flows; the accelerated particle energy E is $\sim q\Phi_0$, where q is the particle charge. This estimate for E derives from the fact that the electrostatic field is dipole-like outside the plasma, giving rise to the potential drop on the order of Φ_0 along B [15].

The process of particle acceleration is transient: The energetic particle outflow from boundary layers of the plasmoid gives rise to plasma current; the resulting force decelerates the plasmoid cross-field motion (see below).

As an example, consider a plasmoid with $h \sim 10$ km infalling at the the free-fall velocity onto the surface of a canonical neutron star of mass $M_* = 1.4M_\odot$, radius $R_* = 10$ km and surface magnetic field $B_* = 5 \times 10^{12}$ G [17]. The accelerating potential $\Phi_0 \sim 10^{21}$ V is then sufficient to accelerate protons to UHECR energies, and heavier nuclei to even higher energies.

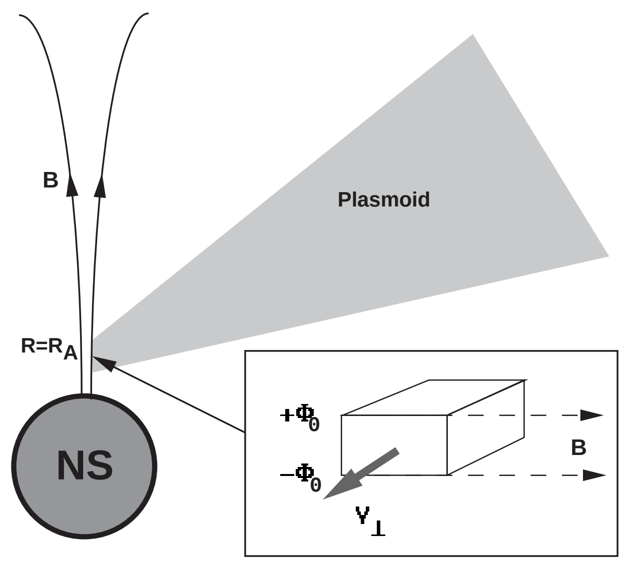

FIG. 1. Schematic of a plasmoid infall on the neutron star.

Several questions arise immediately: Can the infalling plasmoid have the required width? Can accelerated particles escape the neutron star magnetic field? And can the above-described mechanism give rise to the observed energetic particle flux and spectrum? In the remainder of this Letter we address these issues.

Because of their large Larmor radii, the neutron star magnetic field (assumed to be dipolar) would not confine electrons or protons in the considered energy range (although it would confine heavier nuclei). Nevertheless, only a small fraction can escape, due to large radiative losses: the curvature/synchrotron radiation results in a slowing down length that is small when compared to the size of the magnetosphere. Only in regions where the magnetic field curvature is small, viz., near the magnetic axis, are radiative losses not prohibitive.

In order to make our discussion more quantitative, let us consider a specific model. Suppose an iron ($Z = 56$) planetoid (or a planetesimal), with the characteristic mass $M_a \sim 10^{22} - 10^{24}$ g, such as those previously discussed in different contexts [18–22], impacts at the free-fall velocity onto an isolated, slowly rotating neutron star. We assume that the impact is grazing and occurs near the magnetic pole at a large angle ψ ($\gtrsim \pi/6$) to the magnetic field (see Fig. 1).

During the infall the planetoid becomes fragmented and compressed by tidal forces, and ionized, by increased temperature and large motional electric field (the Stark shift becomes comparable to the ionization energy at a distance $R \sim 20 - 100 R_*$ from the neutron star center). At distances larger than the Alfvén radius R_A (i.e., the distance where the ram pressure equals the magnetic pressure and the free-fall velocity equals the Alfvén velocity), the plasmoid motion is expected to be ballistic with the external magnetic field screened from the plasma interior by surface currents [24]. At $R \sim R_A$ the external magnetic field is commonly believed to penetrate the plasma (e.g., [24–26]). The exact mechanism of this penetration (which is likely to involve anomalous resistivity, cf. [27]) is not entirely understood. Nevertheless, we assume that at $R \approx R_A$ the plasma becomes "threaded" by the magnetic field, polarizes and $E \times B$ drifts as described by Chandrasekhar [12] and Schmidt [13].

We adopt the model of Colgate & Petschek [18] to describe the planetoid motion at large distances. Because of a small impact parameter, the planetoid infall is nearly radial. The planetoid becomes fragmented at the distance $R_b = (\rho_0 r_0^2 G M_*/s)^{(1/3}$, where ρ_0 is the density, r_0 is the radius and s is the tensile strength of the planetoid (for an Fe planetoid, $\rho_0 = 8$ g cm^{-3} and $s \sim 10^{10}$ dyne cm^{-2}). The planetoid material initially undergoes a phase of incompressible elongation; it then becomes elongated and compressed at $R < R_i = \kappa R_b$ where $\kappa = (5s/8P_0)^{2/5}$, with P_0 ($\sim 100s$) being the compressive strength. For $R \ll R_i$, $\rho \approx \rho_0 (R_i/R)^{1/2}/4$. Then

$$\frac{R_A}{R_*} = \left(\frac{B_*^2}{\pi \rho_0 V_*^2}\right)^{\frac{2}{9}} \left(\frac{r_0}{R_*}\right)^{-\frac{2}{27}} \left(\frac{2s}{\rho_0 V_*^2 \kappa^3}\right)^{\frac{1}{27}} \quad (1)$$

$$\frac{r_A}{R_*} = \sqrt{5} \left(\frac{B_*^2}{\pi \rho_0 V_*^2}\right)^{\frac{1}{9}} \left(\frac{r_0}{R_*}\right)^{\frac{17}{27}} \left(\frac{2s}{\rho_0 V_*^2 \kappa^{33/20}}\right)^{\frac{5}{27}} \quad (2)$$

where r_A is the planetoid radius at $R = R_A$. We estimate the electrostatic potential drop $\Phi_A = \frac{1}{c} r_A V_A B_A$ where $V_A = V_* (R_*/R_A)^{1/2}$ and $B_A = B_* (R_*/R_A)^3$.

We require that $R_A > R_*$, which implies that

$$B_* > B_{\min} = \left(\frac{r_0}{R_*}\right)^{1/6} \left(\frac{\rho_0 V_*^2 \kappa^3}{2s}\right)^{1/12} \left(\pi \rho_0 V_*^2\right)^{1/2}$$

For an Fe planetoid of mass $\sim 10^{23}$ g and $M_* = 1.4 M_\odot$, $B_{\min} \approx 10^{12}$ G; we therefore limit our attention to $B_* \sim 10^{12} - 10^{14}$ G. For such magnetic fields, $r_A \sim 3 - 8$ km, $R_A \sim 10 - 90$ km and the susceptibility $\chi_\perp(R_A) \sim 3 - 30$.

In order to evaluate the effect of radiative losses on the energy of particles escaping the neutron star magnetosphere we solve the equation of motion of a charged particle including the radiation reaction force:

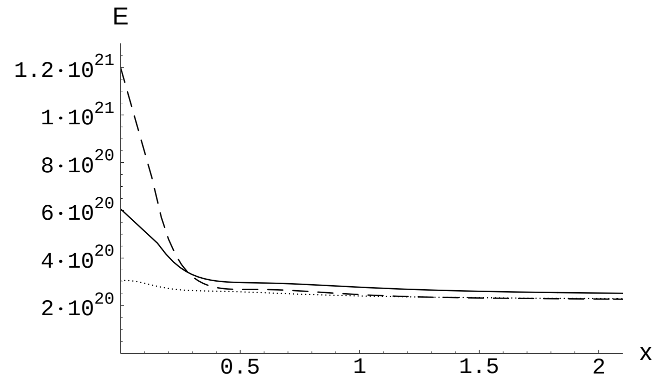

FIG. 2. Particle energy as the function of distance $x = (R - R_0)/R_0$ for $B_* = 5 \times 10^{12}$ G, $\theta = 10^{-4}$ and $M_a = 10^{22}$ g (dotted line), $M_a = 10^{23}$ g (solid line) and $M_a = 10^{24}$ g (dashed line).

$$\frac{d\mathbf{u}}{d\tau} = \mathbf{u} \times \mathbf{\Omega} + \mathbf{F} \quad (3)$$

where τ is the proper time, $\mathbf{u}$ is the spatial component of the four-velocity, $\mathbf{\Omega} = q\mathbf{B}/mc$ is the (vector) cyclotron frequency, m is the particle mass and $\mathbf{F}$ is the radiation reaction force [28]; in the ultra-relativistic limit considered here, $\mathbf{F} = -\lambda (\mathbf{\Omega} \times \mathbf{u})^2 \mathbf{u}$ with $\lambda = 2q^2/3mc^5$.

We solve the above equation for a dipole magnetic field subject to the initial value conditions corresponding to the motion of a particle injected at $t = 0$ in parallel with the magnetic field, with energy $E_0 = Ze\Phi_A$, at the radial distance $R_0 = R_A$ and at an angle θ with respect to the magnetic axis. As we see in Fig. 2, a particle is decelerated by the radiation reaction and emerges from the magnetosphere with a fraction of its injection energy; the emerging particle energy E depends weakly on the planetoid mass. On the other hand, E is sensitive to the injection angle θ (see Fig. 3). For $B_* \sim 10^{12} - 10^{14}$ G, $E > 10^{19}$ eV for angles less than $2 - 6 \times 10^{-3}$.

We can now compute the energy spectrum of particles emerging from the magnetosphere for a *single* impact event. The *average* number of particles per planetoid impact on the neutron star (including impacts not in the vicinity of the magnetic axis), emitted within solid angle $d\Omega$ about the magnetic axis is $dN = N_0 d\Omega/2\pi$, where N_0 is the total number of accelerated particles per event (see below); we have exploited here the fact that the planetoid angular extent is much larger than that of the region in which particles with energies of interest ($\gtrsim 10^{19}$ eV) are generated. Thus the differential energy spectrum is given by $dN/dE = N_0 \sin\theta(E) d\theta/dE$.

In Fig. 4 we show the differential spectrum for magnetic fields $B_* \sim 10^{12} - 10^{14}$ G in the energy range

$10^{19} - 10^{20.5}$ eV observed with AGASA. The spectrum can be well approximated by the power law $dN/dE \sim E^{-\nu}$, with $\nu = 3.03$, 2.95 and 2.89 for $B = 10^{12}$, 10^{13} and 10^{14} G, respectively (the cut-off at $E = 2 \times 10^{20}$ eV for $B = 10^{14}$ G corresponds to $\theta = 0$, cf. Fig. 3). Within the error bars, this agrees with the power spectrum observed with AGASA [29] ($\nu = 2.78^{+0.25}_{-0.33}$) and in the Akeno experiment [30] ($\nu = 2.8 \pm 0.3$). Note, however, that the spectrum of particles emerging from the neutron star magnetosphere might differ from the spectrum observed on Earth, depending on the UHECR confinement characteristics in the galactic magnetic field (see later).

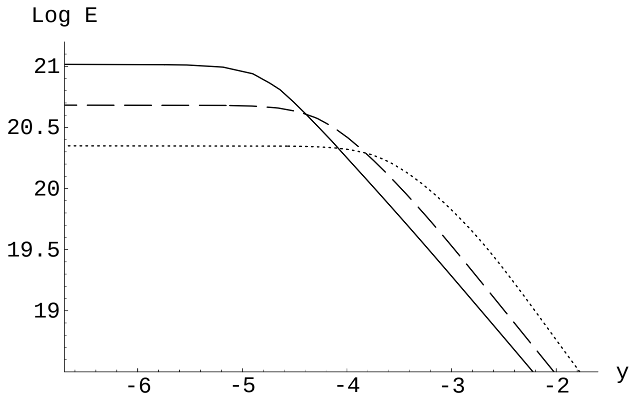

FIG. 3. Energy E (in eV) of a particle emerging from the magnetosphere as the function of the injection angle ($y \equiv \log \theta$), for $B_* = 10^{12}$ G (solid line), $B_* = 10^{13}$ G (dashed line) and $B_* = 10^{14}$ G (dotted line).

Let us now consider whether the mechanism discussed in the present Letter can give rise to the observed flux of the highest energy cosmic rays. AGASA observations [29] imply that the number density of particles with energies exceeding 10^{19} eV is $n_{obs} \approx 6 \times 10^{-29}$ cm^{-3}. In order to determine the density of UHECRs predicted by our model let us first find the number of particles with energies greater than E produced in each impact event,

$$N(E) = \int_E^\infty dE' \frac{dN}{dE'} = N_0[1 - \cos\theta(E)]$$

In order to find N_0 we first note that for the considered range of parameters, the plasmoid is decelerated by the $\mathbf{J} \times \mathbf{B}$ force before reaching the neutron star surface. This can be seen as follows: the field-aligned current carried by the energetic particle outflow from boundary layers of the plasmoid has surface density σc where surface charge density $\sigma = V_\perp B/4\pi c$. Because of charge conservation, this current equals the plasma current, so that the plasma current density $J = 2\sigma c/w$, where w is the plasmoid width along the magnetic field. Because of the $\mathbf{J} \times \mathbf{B}$ force the plasma is decelerated on the time scale $\tau_d = \rho V_\perp c/JB = 2\pi\rho c w/B^2$. In the vicinity of $R = R_A$, $\tau_d = c r_A/V_A^2$. For characteristic fields and planetoid masses considered in this Letter, the deceleration time is shorter than the free-fall time; thus the cross-field motion of the plasmoid is halted before it impacts the neutron star surface. Integrating the momentum balance equation (with gravitational and pressure forces neglected as small compared to the $\mathbf{J} \times \mathbf{B}$ force at $R < R_A$)

$$\rho \frac{d\mathbf{V}}{dt} = \frac{1}{c} \mathbf{J} \times \mathbf{B}$$

over time and plasma volume one then finds that

$$W_A = Q_0 \Phi_A$$

where $W_A = M_a V_A^2/2$ is the planetoid kinetic energy at $R = R_A$ and $Q_0 = ZeN_0$ is the total charge carried by accelerated ions. Integrating the spectrum (Fig. 3), we find that $N(E = 10^{19}\ \text{eV}) \sim 0.2 - 1.4 \times 10^{29}$ for $B_* \sim 10^{12} - 10^{14}$ G.

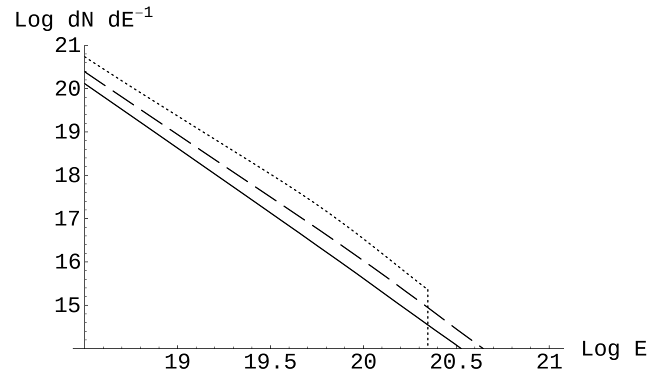

FIG. 4. Differential energy spectrum for $B_* = 10^{12}$ G (solid line), $B_* = 10^{13}$ G (dashed line) and $B_* = 10^{14}$ G (dotted line).

In order to estimate the flux of UHECRs observed at Earth, we need to know the number of neutron stars, the rate of impact events, and the confinement time of cosmic rays in the Galaxy; none of these are known with any degree of certainty.

For the characteristic strength of the galactic magnetic field ($\sim 3\ \mu$G) the Larmor radius of Fe nuclei with $E \lesssim 10^{20}$ eV is $\lesssim 1$ kpc; thus such particles are confined by the galactic magnetic field. The density of UHECRs with energies greater than E is therefore $n = \alpha \tau_c N(E) n_{NS}$, where n_{NS} is the number density of neutron stars, α is the impact rate, and τ_c is the confinement time of UHECRs in the galactic magnetic field. The density of neutron stars is estimated [17] to be $n_{NS} \sim 2 \cdot 10^{-3}$ pc^{-3}. Thus $n = n_{\rm obs}$ requires $\alpha\tau_c \sim 6 - 40$.

The upper bound on the confinement time is given by the UHECR decay time ($\sim 10^{16}$ s) due to photodisintegration by the infrared background radiation [31,32]; the lower bound is set by cross-field particle drifts due to the galactic magnetic field inhomogeneity and curvature. For example, a 10^{19} eV Fe nucleus in a magnetic field with curvature $R_c \sim 10$ kpc would drift a distance on the order

of R_c in 10^{13} s; however, if the magnetic field is twisted, the confinement time could be significantly longer [37].

The confinement time variation with particle energy determines the difference between the spectrum at the source and the spectrum observed on Earth and depends on the character of the global galactic field as well as on the spectrum of magnetic fluctuations. At lower energies (below the "knee") the diffusion time is believed to be a decreasing function of energy [38] which leads to the observed spectrum that is steeper than the source spectrum. For UHECR energies ($> 10^{19}$ eV), however, unlike at lower energies, the gyroradius ρ_c of an Fe nucleus is larger than the integral length scale L_c ($\sim$ 100 pc) of magnetic fluctuations in the Galaxy [39] which may lead to different confinement characteristics [40].

Assuming now the confinement time $\tau_c \sim 10^{13} - 10^{16}$ s, we find the required impact rate to be one in $10^{6\pm2}$ years per neutron star (higher if the density of strongly magnetized neutron stars is significantly lower than n_{NS}).

It is rather difficult to assess whether this impact event rate is plausible. The rate of solid object impacts on neutron stars has been a subject of widely varying estimates in the past [33–36,20,22] in a different context. We do not attempt to make yet another estimate in the present Letter and only note that the rate required by our model is consistent with some of these previous estimates.

In summary, we have explored the plausibility of the acceleration by the polarization electric field, which arises in plasma resulting from planetoid accretion onto magnetized neutron stars, as the generation mechanism for the cosmic rays with highest observed energies. We found that the source spectrum of particles generated by this mechanism is similar to the observed spectrum; whether the resulting spectrum observed on Earth will retain the same character is, at present, an open question. The calculated particle flux magnitude is plausible, albeit quite uncertain due to uncertainties in the UHECR confinement time, in the planetoid impact rate, and in the number of magnetized neutron stars.

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